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Jetzer, Philippe ; Tortora, Crescenzo

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# Limits on decaying dark energy density models from the CMB temperature-redshift relation

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**Abstract.** We discuss the thermodynamic and dynamical properties of a variable dark energy model with density scaling as  $\rho_x \propto (1+z)^m$ ,  $z$  being the redshift. These models lead to the creation/disruption of matter and radiation, which affect the cosmic evolution of both matter and radiation components in the Universe. In particular, we have studied the temperature-redshift relation of radiation, which has been constrained using a recent collection of cosmic microwave background (CMB) temperature measurements up to  $z \sim 3$ . We find that, within the uncertainties, the model is indistinguishable from a cosmological constant which does not exchange any particles with other components. Future observations, in particular measurements of CMB temperature at large redshift, will allow to give firmer bounds on the effective equation of state parameter  $w_{\text{eff}}$  for such types of dark energy models.

## 1. Introduction

The current standard cosmology model relies on the existence of two unknown dark components, the so called “dark matter” (DM) and “dark energy” (DE), which amount to  $\sim 25\%$  and  $\sim 70\%$  of the total energy budget in the Universe, respectively. According to several observations, the Universe is spatially flat and in an accelerated phase of its expansion [1, 2, 3, 4, 5]. DE, described as a cosmological constant  $\Lambda$  in its simplest form, is modelled by a fluid with a negative pressure, which is a fundamental ingredient to explain the actual accelerated expansion of the Universe.

Several models have been proposed to explain DE [6, 7, 8, 9, 10, 11, 12, 13]. An alternative consists to consider a phenomenological variable DE density with continuous creation/disruption of photons [14, 15, 16, 17, 18, 19] or matter [20, 21]. The DE might decay/grow slowly in the course of the cosmic evolution and thus provide the source/sink term for matter and radiation. Different such models have been discussed and strong constraints come from very accurate measurements of the cosmic microwave background (CMB) radiation and other typical cosmological probe.

CMB radiation is the best evidence for an expanding Universe starting from an initial high density state. Within the Friedmann-Robertson-Walker (FRW) models of the Universe the radiation, after decoupling, expands adiabatically and scales as  $(1+z)$ ,  $z$  being the redshift [22]. If we assume that each component is not conserved, contrarily to the standard scenario, then depending on the decay mechanism of the DE, the created photons could lead to distortions in the Planck spectrum of the CMB, and change the evolution of its temperature. The chance to appreciate the deviation from the standard temperature evolution is given by the increasingly number of recent works collecting observations of CMB temperature both at low [23, 24] and higher redshifts [25].

Following the theoretical lines of [14, 15, 16, 17], in Jetzer et al. [18, 19] we have discussed a variable DE model  $\Lambda(z) \propto (1+z)^m$  decaying into photons and DM particles. In particular we studied thermodynamical aspects in the case of a continuous photon creation, which implies a modified temperature redshift relation for the CMB. We have tested the predicted temperature evolution of radiation with some recent data on the CMB at higher redshift from both Sunyaev-Zel'dovich (SZ) effect and high-redshift QSO absorption lines as well as with an updated collection of data from different kind of observations, like distance moduli of Supernovae Ia, observations of the CMB anisotropy and the large-scale structure, together with observational Hubble parameter estimations [19]. A similar approach, however without considering photon creation and thus a modification in the CMB temperature evolution, has been discussed by Ma [21].

## 2. Theoretical framework

We assume a cosmological framework based on the usual Robertson-Walker (RW) metric element [22] and that the Universe contains three different components: a) a matter (both baryons and DM) fluid, with equation of state  $p_m = 0$  (since  $p_m \ll \rho_m$ ), b) a generalised fluid with pressure  $p_\gamma = (\gamma - 1)\rho_\gamma$ , where  $\gamma$  is a free parameter, which is set to 4/3 for a properly said radiation fluid, and c) a DE,  $x$  component, with pressure  $p_x$  and density  $\rho_x$ . The equation of state for the  $x$  component could assume a very general expression, but we limit ourselves to consider the simple linear relation  $p_x = w_x \rho_x$ . We will set any 'bare' cosmological constant  $\Lambda_0$  equal to 0 [18]. With these components we get for the Einstein field equations [22]

$$8\pi G \rho_{tot} = 3 \frac{\dot{R}^2}{R^2} + 3 \frac{k}{R^2}, \quad (1)$$

$$8\pi G p_{tot} = -2 \frac{\ddot{R}}{R} - \frac{\dot{R}^2}{R^2} - \frac{k}{R^2}, \quad (2)$$

where  $p_{tot} = p_\gamma + p_x$  and  $\rho_{tot} = \rho_m + \rho_\gamma + \rho_x$  are the total pressure and density,  $R$  is the scale factor,  $k = 0, \pm 1$  is the curvature parameter and a dot means time derivative. Furthermore, we will assume that there is no curvature, thus  $k = 0$  [3].

In the following we will adopt  $w_x = -1$ , however since we are assuming that the vacuum decays into radiation and massive particles, the effective equation of state  $w_{eff}$ , which is the measured quantity, can differ from  $-1$ .

Following [18, 19] (but see also [14, 16]) we set here the energy conservation equation for the different fluids. The fluid as whole, verifies the Bianchi identity, which means that energy and momentum are locally conserved:  $\nabla_\mu T^{\mu\nu} = 0$ , with  $T^{\mu\nu}$  the stress-energy tensor

$$T^{\mu\nu} = (\rho_{tot} + p_{tot})u^\mu u^\nu - p_{tot}g^{\mu\nu} \quad (3)$$

and  $u^\mu$  being the four velocity. After easy calculations, this conservation law reads

$$\dot{\rho}_{tot} + 3(\rho_{tot} + p_{tot})H = 0, \quad (4)$$

where  $H = \dot{R}/R$  is the Hubble parameter. In the standard approach, each component is conserved, thus Eq. (4) holds for all the fluid components, but here we will suppose that each component will exchange energy with each other. In particular, for matter, radiation and DE we impose the following relations (see also [26])

$$\dot{\rho}_m + 3\rho_m H = (1 - \epsilon) C_x, \quad (5)$$

$$\dot{\rho}_\gamma + 3\gamma\rho_\gamma H = \epsilon C_x, \quad (6)$$

$$\dot{\rho}_x + 3(p_x + \rho_x)H = -C_x, \quad (7)$$

where we assume that both the matter and the radiation fluids exchange energy with the DE as parameterized by  $C_x$  and  $\epsilon$ . In particular,  $C_x$  depends on the DE and, indeed, acts as a source/sink term for the fluids energy. Evidently, if no interaction between the different components exists, then  $C_x$  is null and the standard picture is recovered. Moreover, if  $\epsilon = 0$  ( $\epsilon = 1$ ), then the DE exchanges all the energy with matter (radiation).  $C_x$  can describe different physical situations such as, for instance, a thermogravitational quantum creation theory [15] or a quintessence scalar field cosmology [7]. As we will discuss later on,  $\epsilon$  has to be very small (i.e.,  $\epsilon \ll 1$ ), otherwise the radiation density would become much too big, contrary to present values. As an order of magnitude estimate we expect  $\epsilon \simeq \frac{p_\gamma + \rho_\gamma}{\rho_m}$ , for which indeed  $\epsilon \ll 1$ , since  $p_\gamma + \rho_\gamma \ll \rho_m$ .

Adopting as mentioned above the relation  $p_x = -\rho_x$  and defining  $\rho_x = \Lambda(t)/(8\pi G)$ , from Eq. (7) we obtain

$$C_x = -\frac{\dot{\Lambda}(t)}{8\pi G}. \quad (8)$$

We assume a power law model for the  $\Lambda(t)$  function,  $\Lambda(t) = B(R/R_0)^{-m}$ , or equivalently in terms of redshift,  $\Lambda(z) = B(1+z)^m$ , where  $B$  is a constant. The value of this constant is  $B = 3H_0^2(1 - \Omega_{m0})$ , which can be found using Eq. (1) at the present epoch, assuming that today  $\rho_\gamma \ll \rho_m, \rho_x$ . Thus, the density evolution for the  $x$  component is given by

$$\rho_x/\rho_{crit} = (1 - \Omega_{m0})(1+z)^m, \quad (9)$$

where we have defined as  $\rho_{crit} = \frac{3H_0}{8\pi G}$  the present critical density of the Universe. If  $m$  is positive, then the DE slowly decreases as a function of the cosmic time, whereas if  $m$  is negative the inverse process happens.

From Eqs. (5) and (6) we derive the evolution laws for matter and radiation,

$$\begin{aligned} \rho_m/\rho_{crit} &= \Omega_{m0}(1+z)^3 \\ &- (1-\epsilon)\frac{m(1-\Omega_{m0})}{m-3}[(1+z)^m - (1+z)^3], \end{aligned} \quad (10)$$

$$\begin{aligned} \rho_\gamma/\rho_{crit} &= \Omega_{\gamma0}(1+z)^{3\gamma} \\ &- \epsilon\frac{m(1-\Omega_{m0})}{m-3\gamma}[(1+z)^m - (1+z)^{3\gamma}], \end{aligned} \quad (11)$$

where  $\Omega_{m0}$  and  $\Omega_{\gamma0}$  are the matter and radiation energy densities at  $z = 0$ , respectively.

Since we are interested in the evolution of the radiation temperature, it is useful to discuss the extreme case when only photons enter in the process (i.e.  $\epsilon = 1$ ). Then for the matter density in Eq. (10) the usual evolution  $\propto (1+z)^3$  holds, while the radiation besides the usual term  $\Omega_{\gamma0}(1+z)^{3\gamma}$  has also a perturbative term depending on  $m$ . Since today <sup>1</sup>  $\Omega_{\gamma0} \sim 5 \times 10^{-5}$  it turns out that  $m$  has to be extremely small  $\lesssim 10^{-4}$ . Therefore, unless  $m$  is extremely small or vanishing, DE has to decay mainly in matter with possibly some photons as well. Thus the condition  $\epsilon \ll 1$  has to hold.

<sup>1</sup> The present radiation density is the only cosmological parameter accurately measured. The radiation density is dominated by the energy in the cosmic microwave background (CMB), and the COBE satellite FIRAS experiment determined its temperature to be  $T = 2.725 \pm 0.001$  K [27], corresponding to  $\Omega_{\gamma0} \sim 5 \times 10^{-5}$ .

### 2.1. Hubble and deceleration parameter

Due to the very small value of  $\Omega_{\gamma 0}$  it follows that the evolution of the Universe is essentially driven by the DE and DM, therefore, from Eq. (1) the following law for the Hubble parameter holds

$$H(z) \simeq \frac{8\pi G}{3}(\rho_m + \rho_x) \quad (12)$$

$$= H_0 \left[ \frac{3(1 - \Omega_{m0})}{3 - m}(1 + z)^m + \frac{(3\Omega_{m0} - m)}{3 - m}(1 + z)^3 \right]^{1/2},$$

which is obviously the same expression found in Ma [21].

Recasting Eq. (7), it is possible to write

$$\dot{\rho}_x + 3H(p_x + \rho_x + \frac{C_x}{3H}) = 0, \quad (13)$$

which shows that the term  $C_x$  contributes to an effective pressure

$$p_{\text{eff}} = p_x + \frac{C_x}{3H} = -\rho_x + \frac{C_x}{3H}. \quad (14)$$

Therefore, we get an equivalent effective DE equation of state  $w_{\text{eff}}$  [21]

$$w_{\text{eff}} = \frac{p_{\text{eff}}}{\rho_x} = \frac{m}{3} - 1. \quad (15)$$

If  $m > 0$  then we have  $w_{\text{eff}} > -1$ , i.e. our model is quintessence-like [6, 7, 12], while we have a phantom-like [9] model when  $m$  is negative and  $w_{\text{eff}} < -1$ . Another interesting quantity is the deceleration parameter, which can be written as

$$q(z) = -\frac{\ddot{R}R}{\dot{R}^2} = \quad (16)$$

$$\frac{(1+z)^3(m - 3\Omega_{m0}) + 3(m-2)(1+z)^m(\Omega_{m0} - 1)}{2(1+z)^3(m - 3\Omega_{m0}) + 6(1+z)^m(\Omega_{m0} - 1)}.$$

Imposing that  $q(z) = 0$ , we can determine the *transition redshift*, i.e. the redshift at which the Universe changed from a deceleration to an acceleration phase, which is given by

$$z_T = \left( \frac{3(2-m)(1-\Omega_{m0})}{3\Omega_{m0}-m} \right)^{\frac{1}{3-m}} - 1. \quad (17)$$

From this result we see that the larger  $m$  is, the earlier the Universe changes from deceleration to acceleration.

### 2.2. Thermodynamical aspects and CMB temperature evolution

We follow here the approach outlined in Lima [14], where he defines a current as  $N^\alpha = nu^\alpha$  with  $n$  being the particle number density of the photons or of the DM particles. Indeed, there is a current for each of these components. Due to the decaying vacuum the current satisfies the following balance equation (one for each component)

$$\dot{n}_i + 3n_i H = \psi_i, \quad (18)$$

with  $i = \gamma$  or  $DM$  ( $\gamma$  for the photons and  $DM$  for the dark matter) and  $\psi_i$  is the corresponding particle source. For decaying vacuum models  $\psi_\gamma + \psi_{DM}$  is positive and related to the rate of change of  $\rho_x$ . We can also define an entropy current of the form

$$S^\alpha = \sum_i n_i \sigma_i u_i^\alpha, \quad (19)$$

where  $\sigma_i$  is the specific entropy per particle (photons, DM and in principle also DE). If the DE  $\rho_x$  is constant the above entropy current is conserved. The existence of a non equilibrium decay process of the vacuum implies  $S_{;\alpha}^\alpha \geq 0$ , thus an increase of the entropy as a consequence of the second law of thermodynamics. In principle the second law should be applied to the system as a whole, thus including the vacuum component [26]. Assuming that the vacuum is like a condensate with zero chemical potential  $\mu_{vac}$  it follows from Euler's relation

$$\mu_{vac} = \frac{\rho_x + p_x}{n} - T\sigma_{vac}, \quad (20)$$

provided  $w_x = -1$ , that  $\sigma_{vac} = 0$  and thus its contribution to the entropy current vanishes. Given our assumptions the vacuum plays the role of a condensate carrying no entropy.

For instance, for quintessence models in the limit where the scalar field does not depend on time, and thus its time derivative vanishes, one gets  $w_x = -1$ . In the later stages of the Universe the time dependence is possibly very weak so that  $w_x = -1$  holds up to small corrections.

The equation for the particle number density of radiation component is given by Eq. (18) with  $i = \gamma$ . Using Gibbs law and well-known thermodynamic identities, following the derivation given in the paper by Lima et al. [16], one gets (see also [14, 18])

$$\frac{\dot{T}}{T} = \left( \frac{\partial p_\gamma}{\partial \rho_\gamma} \right)_n \frac{\dot{n}_\gamma}{n_\gamma} - \frac{\psi_\gamma}{n_\gamma T \left( \frac{\partial \rho_\gamma}{\partial T} \right)_n} \left[ p_\gamma + \rho_\gamma - \frac{n_\gamma \epsilon C_x}{\psi_\gamma} \right]. \quad (21)$$

To get a black-body spectrum the second term in brackets in Eq. (21) has to vanish, thus

$$\epsilon C_x = \frac{\psi_\gamma}{n_\gamma} [p_\gamma + \rho_\gamma]. \quad (22)$$

Thus, Eq. (21) becomes

$$\frac{\dot{T}}{T} = \left( \frac{\partial p_\gamma}{\partial \rho_\gamma} \right)_n \frac{\dot{n}_\gamma}{n_\gamma}. \quad (23)$$

With  $\left( \frac{\partial p_\gamma}{\partial \rho_\gamma} \right)_n = (\gamma - 1)$  one obtains

$$\frac{\dot{T}}{T} = (\gamma - 1) \frac{\dot{n}_\gamma}{n_\gamma}. \quad (24)$$

Using the equation for the particle number conservation Eq. (18) into Eq. (24) leads to

$$\frac{\dot{T}}{T} = (\gamma - 1) \left[ \frac{\psi_\gamma}{n_\gamma} - 3H \right]. \quad (25)$$

With Eqs. (22) and (25) we get

$$\frac{\dot{T}}{T} = (\gamma - 1) \left[ -\frac{\epsilon \dot{\Lambda}}{8\pi G(p_\gamma + \rho_\gamma)} - 3H \right]. \quad (26)$$

Now, following the previous discussion on  $\epsilon$  and aiming to be very general, we set  $\epsilon = \frac{\rho_\gamma + p_\gamma}{\rho_m} \tilde{\epsilon}$ , where  $\tilde{\epsilon}$  is a new parameter and insert it into Eq. (26). Taking the sum of Eqs. (1) and (2) we find

$$8\pi G(\rho_{tot} + p_{tot}) \simeq 8\pi G\rho_m = 2\frac{\dot{R}^2}{R^2} - 2\frac{\ddot{R}}{R} = -2\dot{H} . \quad (27)$$

Finally, we obtain the expression

$$\frac{\dot{T}}{T} = (\gamma - 1) \left[ \frac{\dot{\Lambda}\tilde{\epsilon}}{2\dot{H}} - 3H \right] , \quad (28)$$

which we can integrate

$$\int_{t_1}^{t_0} \frac{\dot{T}}{T} dt = (\gamma - 1) \int_{t_1}^{t_0} \left[ \frac{\dot{\Lambda}\tilde{\epsilon}}{2\dot{H}} - 3H \right] dt , \quad (29)$$

where  $t_0$  denotes the present time and  $t_1$  some far instant in the past. Indeed, if  $\dot{\Lambda}$  vanishes and  $\gamma = 4/3$  one gets the usual dependence  $T(t) = \frac{R(t_1)T(t_1)}{R(t)}$  for a radiation fluid. To carry out the integration of the first term on the right hand side it is useful to perform a change of variable from  $t$  to  $z$  and accordingly  $\frac{dt}{dz} = \frac{-1}{H(1+z)}$ . This way we get (with  $z_1$  corresponding to the time  $t_1$  and  $z_0 = 0$  corresponding to  $t_0$  present time)

$$\begin{aligned} \ln \frac{T(z=0)}{T(z_1)} + 3(\gamma - 1) \ln \frac{R(z=0)}{R(z_1)} = \\ \frac{(\gamma - 1)}{2} \int_0^{z_1} \frac{\Lambda'\tilde{\epsilon}}{H'H(1+z)} dz , \end{aligned} \quad (30)$$

where  $'$  denotes derivative with respect to  $z$ .

As next we insert  $H(z)$  and its derivative as taken from Eq. (12) into Eq. (30) and integrate it, to get (setting  $z_1 = z$ )

$$T(z) = T_0 \left( \frac{R_0}{R(z)} \right)^{3(\gamma-1)} \exp \left( \frac{B(1-\gamma)\tilde{\epsilon}}{3H_0^2(\Omega_{m0}-1)} A \right) , \quad (31)$$

where

$$\begin{aligned} A = \ln((m - 3\Omega_{m0}) + m(1+z)^{m-3}(\Omega_{m0}-1)) \\ - \ln((m - 3)\Omega_{m0}) . \end{aligned} \quad (32)$$

We can also write Eq. (31) as

$$\begin{aligned} T(z) = T_0(1+z)^{3(\gamma-1)} \\ \times \left( \frac{(m - 3\Omega_{m0}) + m(1+z)^{m-3}(\Omega_{m0}-1)}{(m - 3)\Omega_{m0}} \right)^{\tilde{\epsilon}(\gamma-1)} . \end{aligned} \quad (33)$$

We inserted in the exponent of Eq. (31) the explicit form of  $B$ , thus getting as exponent in the above Eq.  $\tilde{\epsilon}(\gamma - 1)$ . Hereafter, we will set  $\tilde{\epsilon} = 1$ . Clearly  $\tilde{\epsilon}$  and  $m$  are not independent, we checked using the temperature redshift data that if  $\tilde{\epsilon}$  is bigger ( $\sim 10$  or more), then  $m$  has to be extremely small consistently with what mentioned above (as it would lead to a too high production of photons in the DE decay). On the other hand, if  $\tilde{\epsilon}$  gets smaller (e.g.,  $\sim 0.1$ )  $m$  gets bigger ( $\sim 0.2$ ) and accordingly  $w_{eff}$ , moreover  $m$  would be poorly constrained, since the

uncertainties would then be very high. But from the other data (without the the temperature ones) there are already stringent limits on  $m$  and thus this way one could get lower limits on  $\tilde{\epsilon}$ , under the assumption that DE decays also in photons.

Notice that for  $z = 0$  we have  $T(0) = T_0$ , whereas for  $m = 0$  the expression in the parenthesis is equal to 1 and thus  $T(z) = T_0(1+z)^{3(\gamma-1)}$ , which for the canonical value of  $\gamma = 4/3$  reduces to the standard expression.

### 3. Results

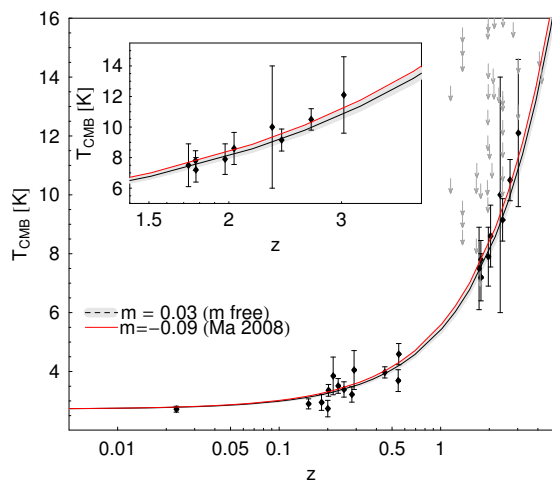
To test the temperature evolution for the radiation component, we rely on the CMB temperatures derived from the absorption lines of high redshift systems and the ones from SZ effect in clusters of galaxies (we will collectively quote as  $T_{CMB}$ , hereafter). At high redshift the CMB temperature is recovered from the excitation of interstellar atomic or molecular species that have transition energies in the sub-millimetre range and can be excited by CMB photons. When the relative population of the different energy levels are in radiative equilibrium with the CMB radiation, the excitation temperature of the species equals that of the black-body radiation at that redshift, providing one of the best tools for determining the black-body temperature of the CMB in the distant Universe [28, 29, 30, 31, 32, 33, 34, 35]. We used a sample of 9 QSO absorption measurements [19, 25]. In summary we have 4 data points from the analysis of the fine structure of atomic carbon (AC) and 5 measurements based on the rotational excitation of CO molecules (CO) [25].

At lower redshift we use the measurements from the SZ effect. Thus, spectral measurements of galaxy clusters at different frequency bands yield independent intensity ratios for each cluster. The combinations of these measured ratios permit to extract the cosmic microwave background radiation (see Fabbri et al. [36]). We relied on the data compilation in Luzzi et al. [24], which have analyzed the results of multifrequency SZ measurements toward several clusters from 5 telescopes (BIMA, OVRO, SUZI II, SCUBA and MITO).

We derive the observed  $T_{CMB}$  from the theoretical expression  $T_{th}$ , which we have derived in Eq. (33). We set  $T_0 = 2.725$  K, which is quite well determined in the literature [37], and the matter density  $\Omega_{m0} = 0.273$  to the value inferred in Komatsu et al. [38]. If we take  $\gamma = 4/3$ , then we find  $m = 0.03^{+0.08}_{-0.09}$  [19], which is lower than the estimated value of  $m = 0.09 \pm 0.10$  in [18], but fully consistent within uncertainties, and also pretty consistent with  $m = 0$ . In Fig. 1 the temperature measurements (together with some upper limits) are shown, and our best fitted result is plotted and compared with the  $m = -0.09$  result in Ma [21]. The value we have found corresponds to an effective equation of state  $w_{\text{eff}} = -0.99 \pm 0.03$ , consistent with  $w_{\text{eff}} = -1$ , and the transition redshift is  $z_T = 0.78 \pm 0.08$ . In order to check the impact of redshift distribution we separate the data in two redshift bins with  $z < 0.6$  (SZ data only) and  $z \geq 0.6$  (QSO absorption lines only), finding the best fitted values  $m = 0.12^{+0.12}_{-0.13}$  and  $-0.05^{+0.12}_{-0.14}$ , respectively. Although the uncertainties are very high, it seems to emerge a mild trend with lower redshift data preferring a DE decaying into matter and radiation, while data at  $z > 0.6$  point to an opposite behavior. These results could be interpreted in a different way, in fact, the differences found could be due not to the different redshift coverage, but to some particular biases in the two kind of observations, SZ vs QSO absorption lines. An indication of this suggestion comes if we divide the sample in three subsamples: 1) SZ data, 2) the data from the analysis of the fine structure of atomic carbon (AC), and 3) the measurements based on the rotational excitation of CO molecules in [25]. If we fit the model to the combined SZ+AC and SZ+CO samples we find  $m = 0.11^{+0.11}_{-0.11}$  and  $m = 0.03^{+0.09}_{-0.10}$ , respectively. Because of the larger measurement errors, the AC data affect very little the results when only SZ measurements are adopted, and the result for the second sample shows that SZ and CO data mainly constrain  $m$ , and AC simply gives a tiny reduction on the errors.

Adopting a constant value for the ratio  $\psi_\gamma/3n_\gamma H = \beta$  Lima et al. [16] have found the simple





**Figure 1.** Cosmic microwave background temperature as a function of the redshift. The black points with bars are the full collection of measurements from Luzzi et al. [24] and Noterdaeme et al. [25]. The gray arrows represent the upper limits derived from the analysis of atomic carbon (see [25] for details). The black line is the best fit result ( $m = 0.03$ ), while the gray region is the  $1\sigma$  uncertainty. The red line is the best fit recovered from Ma [21] (see the legend). The inset panel shows a magnified vision of the higher redshift region of the plot.

relation,

$$T(z) = T_0(1+z)^{1-\beta}. \quad (34)$$

If we fit all the sample we find  $\beta = -0.002 \pm 0.03$ , while for the low and high redshift subsamples we have  $\beta = 0.06 \pm 0.08$  and  $-0.01 \pm 0.03$ , respectively. These results are qualitatively consistent with what found in Noterdaeme et al. [25], and points to a similar trend as the one discussed above.

If  $\gamma$  is left free to change, for the best fitted value it turns out that  $\gamma > 4/3$  and  $m$  is systematically more positive. We find  $\gamma = 1.35^{+0.03}_{-0.03}$  and  $m = 0.25^{+0.23}_{-0.17}$ , which corresponds to an effective equation of state  $w_{\text{eff}} = -0.92 \pm 0.07$  and the transition redshift is  $z_T = 1.1 \pm 0.6$ . When the two subsamples are adopted, wide confidence contours are found, particularly for the  $z \geq 0.6$  sample, for which the contours at very low  $m$  are not closed. We obtain  $\gamma = 1.3^{+0.2}_{-0.1}$  and  $1.26^{+0.01}_{-0.01}$ , while  $m = 0.8^{+0.1}_{-0.3}$  and  $0.6^{+0.1}_{-1.0}$ , respectively for the low and high- $z$  samples.

We notice that if DE does not decay into radiation (corresponding to  $\epsilon = \tilde{\epsilon} = 0$  and thus  $m$  is no longer constrained) then the CMB temperature will scale in the standard way. Clearly, this would imply that if DE decays, this has to be into DM only. On the other hand a deviation of the CMB temperature from the standard scaling could be interpreted as DE decaying also into radiation. In which case with Eq. (33) one can determine either  $m$  or  $\tilde{\epsilon}$  and thus get some insights on the decay mode of DE into radiation. Future data on the CMB temperature will allow to shed light on this important issue.

In [19] we have combined the CMB temperature data with a set of other independent measurements such as Supernovae Ia distance moduli, CMB anisotropy, baryon acoustic oscillation and observational data for the Hubble parameter. When combining all the data sets we found values for  $m$ , and thus  $w_{\text{eff}}$  consistent with the ones obtained when using only CMB temperature data.

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